

THE CHARMED STRANGE MESON SYSTEM

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Abstract

Motivated by the recent surprising discovery of two new meson states with $c\bar{s}$ quark content but unexpectedly low masses and narrow total decay widths, we work out, in a nonrelativistic potential-model approach developed already some two decades ago, the predictions for the energy levels of the corresponding charm–anti-strange quark bound states. In spite of the fact that this simple quark model reproduces the mass spectrum of the previously observed hadrons remarkably well, we are led to the conclusion that, without considerable modifications, both the new states do not fit into this framework.

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1 Introduction

Recently, evidence for the existence of two new mesonic states with charm and strange quark content has been presented by three experiments [1–13]. Table 1 summarizes the experimental measurements of the masses of the two new states as well as our weighted averages of the meson masses. The uncertainties of the averages have been obtained by adding in quadrature the statistical and systematic errors reported by the experiments.

At first, the BABAR Collaboration [1], running the BABAR detector at the PEP-II asymmetric-energy e^+e^- storage ring, announced the observation of a charmed strange meson called $D_{sJ}^*(2317)^\pm$. This new state manifested as narrow resonance with a mass of $M(D_{sJ}^*(2317)^\pm) = (2316.8 \pm 0.4 \pm 3.0)$ MeV in the $D_s^\pm \pi^0$ invariant-mass distribution.

The CLEO Collaboration [2,3,12] provided the confirmation of the existence of the $D_{sJ}^*(2317)^\pm$ meson from data collected with the CLEO II detector in symmetric e^+e^- collisions at the Cornell Electron Storage Ring. Furthermore, the CLEO Collaboration [2,3,12] reported the observation of a charmed, strange meson called $D_{sJ}(2463)^\pm$. This new state manifested itself as narrow resonance in the $D_s^\pm \pi^0$ invariant-mass spectrum. Interestingly, the CLEO experiment [3,12] found, within errors, identical results for the mass differences (with statistical and systematic errors) between these two $D_{sJ}^*(2317)^\pm$ and $D_{sJ}(2463)^\pm$ states and the well-established [14] D_s^\pm and $D_s^{*\pm}$ mesons, respectively:

$$\begin{aligned} M(D_{sJ}^*(2317)^\pm) - M(D_s^\pm) &= (350.0 \pm 1.2 \pm 1.0) \text{ MeV} , \\ M(D_{sJ}(2463)^\pm) - M(D_s^{*\pm}) &= (351.2 \pm 1.7 \pm 1.0) \text{ MeV} . \end{aligned}$$

With the most recent Particle Data Group averages [14] $M(D_s^\pm) = (1968.1 \pm 0.5)$ MeV, for the mass of the D_s^\pm meson, and $M(D_s^{*\pm}) = (2111.9 \pm 0.7)$ MeV, for the mass of the $D_s^{*\pm}$ meson, the mass differences translate into the resonance masses listed in Table 1. In return the BABAR experiment confirmed the existence of the $D_{sJ}(2463)^\pm$ meson [7].

Finally, both discoveries have been confirmed by the Belle Collaboration, operating the Belle detector at the KEKB asymmetric-energy e^+e^- collider; the Belle experiment identified the new $D_{sJ}^*(2317)^\pm$ and $D_{sJ}(2463)^\pm$ states in the exclusive B meson decays $B \rightarrow D + D_{sJ}^*(2317)$ and $B \rightarrow D + D_{sJ}(2463)$ [5,11], and continuum e^+e^- annihilations [6].

Table 1: Masses of the two new narrow charmed, strange meson states $D_{sJ}^*(2317)^\pm$ and $D_{sJ}(2463)^\pm$ with both statistical and systematic errors, measured by the experiments BABAR [1,7], CLEO [2,3,8,12], and Belle [5,6,9,11], as well as their weighted averages

Experiment	Process	$M(D_{sJ}^*(2317)^\pm)$ [MeV]	$M(D_{sJ}(2463)^\pm)$ [MeV]
BABAR [1,7]	inclusive e^+e^-	$2316.8 \pm 0.4 \pm 3.0$	$2457.0 \pm 1.4 \pm 3.0$
CLEO [3,8,12]	inclusive e^+e^-	$2318.1 \pm 1.2 \pm 1.0$	$2463.1 \pm 1.7 \pm 1.0$
Belle [5,11]	B decays	$2319.8 \pm 2.1 \pm 2.0$	$2459.2 \pm 1.6 \pm 2.0$
Belle [6]	inclusive e^+e^-	$2317.2 \pm 0.5 \pm 0.9$	$2456.5 \pm 1.3 \pm 1.1$
Weighted Averages		2317.6 ± 0.8	2459.0 ± 1.1

Until now, for the $D_{sJ}^*(2317)^\pm$ state, only the decay $D_{sJ}^*(2317)^+ \rightarrow D_s^+ + \pi^0$ and, for the $D_{sJ}(2463)^\pm$ meson, the decays $D_{sJ}(2463)^+ \rightarrow D_s^{*+} + \pi^0$ and $D_{sJ}(2463)^+ \rightarrow D_s^+ + \gamma$ (as well as, of course, their charge-conjugate channels) have been observed [1–9, 11, 12]. For both states, the widths of the resonance peaks are consistent with the experimental mass resolution. This indicates that the observed resonance widths are consistent with those expected for states of zero intrinsic or natural width Γ . The CLEO Collaboration reports identical upper limits: $\Gamma(D_{sJ}^*(2317)^\pm) < 7$ MeV for the $D_{sJ}^*(2317)^\pm$ meson and $\Gamma(D_{sJ}(2463)^\pm) < 7$ MeV for the $D_{sJ}(2463)^\pm$ meson, at the 90% confidence level [3, 12].

Concerning the quantum numbers of these two narrow resonances, all experimental findings, such as all the observed decay modes and angular distributions, are consistent with their interpretation as P-wave states with spin-parity assignment $J^P = 0^+$ for the $D_{sJ}^*(2317)^\pm$ meson [1–4, 6–9, 12] and $J^P = 1^+$ for the $D_{sJ}(2463)^\pm$ meson [2–9, 11–13].

The new states do not show the decay patterns favoured by theoretical expectation:

- The mass of the $D_{sJ}^*(2317)^\pm$ meson is 40.8 MeV below the kinematical threshold 2358.4 MeV for its isospin-conserving strong decay $D_{sJ}^*(2317)^+ \rightarrow D^0 + K^+$, and 49.4 MeV below the kinematical threshold 2367.0 MeV for its isospin-conserving strong decay $D_{sJ}^*(2317)^+ \rightarrow D^+ + K^0$ (charge-conjugate decays are understood).
- The mass of the $D_{sJ}(2463)^\pm$ meson lies above the two kinematical thresholds for the decays $D_{sJ}(2463)^+ \rightarrow D^0 + K^+$ and $D_{sJ}(2463)^+ \rightarrow D^+ + K^0$; these processes, however, are in conflict with the conservation of angular momentum and parity.
- The mass of the $D_{sJ}(2463)^\pm$ meson is 41.5 MeV below the kinematical threshold 2500.5 MeV for its isospin-conserving decay $D_{sJ}(2463)^+ \rightarrow D^*(2007)^0 + K^+$, and 48.7 MeV below the kinematical threshold 2507.7 MeV for its isospin-conserving decay $D_{sJ}(2463)^+ \rightarrow D^*(2010)^+ + K^0$ (charge-conjugate decays are understood).

As a consequence of their low masses blocking their isospin-conserving strong decays, the mesons $D_{sJ}^*(2317)^\pm$ and $D_{sJ}(2463)^\pm$ decay via the channels $D_{sJ}^*(2317)^+ \rightarrow D_s^+ + \pi^0$ and $D_{sJ}(2463)^+ \rightarrow D_s^{*+} + \pi^0$, which both violate isospin conservation, and through the radiative decay mode $D_{sJ}(2463)^+ \rightarrow D_s^+ + \gamma$, as well as the respective charge-conjugate modes. This explains the observed narrowness of the total decay widths of these states.

The experimentally measured masses of the two $D_{sJ}^*(2317)^\pm$ and $D_{sJ}(2463)^\pm$ states (summarized in Table 1) are significantly smaller than the corresponding predictions of most, but not all, QCD-inspired quark potential models, which regard these mesons as P-wave bound states $(c\bar{s})$ of a charm quark c and a strange antiquark \bar{s} . Concerning the masses of the $J^P = 0^+$ and the lower $J^P = 1^+$ $(c\bar{s})$ states, typical sets of predictions are obtained, e.g., in the comprehensive potential-model analyses of Refs. [15–17], whereas results at the low-mass end of the spectrum of predictions can be found in Refs. [18, 19]. The failure of quark potential models to explain such low masses of the new resonances prompted a multitude of speculations about the genuine nature of these states [20–58]. Tentative interpretations view the mesons $D_{sJ}^*(2317)^\pm$ and $D_{sJ}(2463)^\pm$ dominantly as $(c\bar{s})$ bound states [20, 22, 24, 26–28, 31, 36, 38, 46, 49, 51], (DK) molecules [21], four-quark bound states [23, 29, 33, 53], or $D\pi$ atoms [25]. Clearly, in general each of these mesons is a linear combination of all Fock states which carry the appropriate quantum numbers.

In this work, we examine whether the new D_s mesons can be accommodated as $(c\bar{s})$ bound states in a phenomenologically appealing nonrelativistic quark potential model. After recalling in Sec. 2 the spectroscopic classification of heavy-light states, we sketch in Sec. 3 the adopted quark model, and present in Sec. 4 our—disappointing—findings.

2 Spectroscopy of heavy-quark light-quark mesons

The spectroscopic classification of bound states $(Q\bar{q})$ composed of a heavy quark Q and a comparatively light antiquark \bar{q} is greatly facilitated by the following observation. In the limit of infinitely large mass of the heavy quark Q , the spin \mathbf{s}_Q of the heavy quark Q and the total angular momentum \mathbf{j} of the light quark q , defined as the sum $\mathbf{j} = \mathbf{l} + \mathbf{s}_q$ of the orbital angular momentum \mathbf{l} and the spin \mathbf{s}_q of the light quark q , are separately conserved. Hence, in this case \mathbf{j} provides a “good” quantum number for classification.

The light-quark angular momentum \mathbf{j} and the spin \mathbf{s}_Q of the heavy quark combine to the total angular momentum $\mathbf{J} = \mathbf{j} + \mathbf{s}_Q$ of the $(Q\bar{q})$ bound state. The corresponding quantum number J fixes the spin of the resulting meson. The parity P of this meson is related to the quantum number ℓ of the orbital angular momentum \mathbf{l} by $P = (-1)^{\ell+1}$.

Thus, in the heavy-quark limit the heavy-light bound states $(Q\bar{q})$ may conveniently be classified in terms of the light-quark total-angular-momentum quantum number j :

S-wave states have orbital angular momentum $\ell = 0$ and, therefore, negative parity, $P = -1$. Spin and orbital angular momentum of the light quark q can couple only to $j^P = \frac{1}{2}^-$. Combining the light-quark total angular momentum $\mathbf{j} = \mathbf{s}_q$ with the spin of the heavy quark Q yields a spin-singlet state with spin-parity assignment $J^P = 0^-$ and a spin-triplet state with spin-parity assignment $J^P = 1^-$.

The pseudoscalar $J^P = 0^-$ state has been identified with the isosinglet D_s^\pm meson (the former F^\pm meson), with a mass of (1968.1 ± 0.5) MeV and a well-established spin-parity assignment $J^P = 0^-$ [14]. The vector $J^P = 1^-$ state is assumed to be identical to the isosinglet $D_{s1}^{*\pm}$ meson, with a mass of (2111.9 ± 0.7) MeV, natural spin-parity, and width and decay modes consistent with the assignment $J^P = 1^-$ [14].

P-wave states have orbital angular momentum $\ell = 1$ and, therefore, positive parity, $P = +1$. Spin and orbital angular momentum of the light quark can couple either to $j^P = \frac{1}{2}^+$ or to $j^P = \frac{3}{2}^+$. Combination with the spin of the heavy quark yields, for $j^P = \frac{1}{2}^+$, two states with spin-parity assignment $J^P = 0^+$ and $J^P = 1^+$, and, for $j^P = \frac{3}{2}^+$, two states with spin-parity assignment $J^P = 1^+$ and $J^P = 2^+$. The two $J^P = 1^+$ states do not have definite charge-conjugation properties; therefore, they can undergo mixing.

The vector $J^P = 1^+$ state belonging to the $j^P = \frac{3}{2}^+$ doublet (with possibly small admixtures of its vector $J^P = 1^+$ counterpart belonging to the $j^P = \frac{1}{2}^+$ doublet) is, in general, assumed to be identical to the isosinglet $D_{s1}(2536)^\pm$ meson, with a mass of $(2535.35 \pm 0.34 \pm 0.5)$ MeV and its spin-parity assignment $J^P = 1^+$ which is strongly favoured but still needs confirmation [14].

The spin-triplet tensor $J^P = 2^+$ state is identified with the isosinglet $D_{sJ}(2573)^\pm$ meson, with a mass of (2572.4 ± 1.5) MeV, natural spin-parity, and width and decay modes consistent with the assignment $J^P = 2^+$ [14].

Therefore, the present phenomenological status of the D_s system can be summarized in the following way: The two S-wave (ground) states are experimentally well established. Good candidates for two of the four P-wave states predicted by theory for bound states of two spin- $\frac{1}{2}$ fermions have been observed. The open question is: May the $D_{sJ}^*(2317)^\pm$ and $D_{sJ}(2463)^\pm$ states be interpreted as one or the other of the two missing $(c\bar{s})$ states?

3 Nonrelativistic quark–antiquark potential model

The present analysis of the charmed strange meson system is based on a nonrelativistic quark model improved by the lowest-order relativistic corrections. For recent reviews of the phenomenological description of hadrons as bound states of quarks within the more intuitive framework of (both nonrelativistic and semirelativistic) potential models see, for instance, Refs. [59,60]. Since we intend to describe bound states composed of both a heavy quark and a light quark, we have to make use of a framework which is capable to reproduce the properties not only of heavy quarkonia but also of light mesons [61–63].

The nonrelativistic interquark potential $V(r)$ characterizing the phenomenological model employed in the present analysis is of a “funnel-like” Coulomb-plus-linear shape, improved by an intermediate exponential damping [61]. This quark model provides an excellent description of both the observed [14] meson and baryon mass spectra [62,63]. Apart from the strong fine-structure constant α_s and the masses of the quarks, this model involves six constant parameters, a, b, c, d, V_0, R_1 , entering in the potential $V(r)$:

$$\begin{aligned} V(r) &= -\frac{4}{3} \frac{\alpha_s}{r} - b \exp\left(-\frac{r}{c}\right) + d + V_0 \quad \text{for } r \leq R_1, \\ V(r) &= a r - b \exp\left(-\frac{r}{c}\right) + V_0 \quad \text{for } r \geq R_1. \end{aligned}$$

From the smoothness of $V(r)$ at $r = R_1$, the potential parameters R_1 and d are related to the strong fine-structure constant α_s and the slope a of the linear term according to

$$R_1 = \sqrt{\frac{4\alpha_s}{3a}}, \quad d = 4\sqrt{\frac{\alpha_s a}{3}}.$$

Starting from the Bethe–Salpeter formalism for the description of bound states within relativistic quantum field theory, the static potential $V(r)$ in the Schrödinger equation may be derived from the interaction kernel entering in the Bethe–Salpeter equation by Fourier transformation. However, for a confining potential, such as the linear potential $V(r) = a r$, the corresponding Bethe–Salpeter kernel has to be a highly singular object. The necessary regularization of the kernel introduces a constant in the static potential. In order to compensate the (infrared) divergence of the interaction kernel, in the case of the linear potential $V(r) = a r$ this (negative) constant V_0 has to involve the slope a of the linear potential, and the Euler–Mascheroni constant $\gamma = 0.577\,215\,664\,901\dots$ [64]:

$$V_0 = -2\sqrt{a} \exp\left[-\left(\gamma - \frac{1}{2}\right)\right] \simeq -2\sqrt{a}.$$

The above relations reduce the number of parameters in the potential to four: α_s, a, b, c . With respect to the Lorentz structure of the quark–antiquark interaction, the potential $V(r)$ is split into a Lorentz-scalar part $V_S = a r$ and a Lorentz-vector part $V_V = V - V_S$.

The generalized Breit–Fermi Hamiltonian contains all relativistic corrections to the (nonrelativistic) Schrödinger Hamiltonian up to and including the second order in the inverse masses of the bound-state constituents. This lowest-order relativistic correction discriminates between different Lorentz structures of the quark–antiquark interaction. The spin-dependent relativistic corrections consist, for a Lorentz-vector spin structure, of a spin–orbit, a spin–spin, and a tensor term, and, for a Lorentz-scalar spin structure, of only a spin–orbit term. The spin–spin term, however, involves the Laplacian $\Delta V_V(r)$

of the vector potential $V_V(r)$; for any Coulomb-like contribution to the vector potential $V_V(r)$, i.e., a term in $V_V(r)$ proportional to $1/r$, this yields a highly singular expression:

$$\Delta \frac{1}{r} = -4\pi \delta^{(3)}(\mathbf{r}) .$$

Similarly, for a Coulomb-like contribution $V(r) \propto 1/r$ to the nonrelativistic interaction potential $V(r)$ the spin–orbit and tensor terms introduce singularities of the form $1/r^3$. This makes the Breit–Fermi Hamiltonian an operator which is unbounded from below. However, for heavy bound-state constituents, all relativistic corrections may be treated as perturbations of the well-defined nonrelativistic Schrödinger Hamiltonian operator.

The mass difference between the pseudoscalar π and the vector ρ meson, generated by the spin–spin interaction, is of the same order of magnitude as the involved meson masses. This indicates that, for any description of mesons containing light quarks, the inclusion of at least the spin–spin term in the nonperturbative treatment is mandatory. Within the nonrelativistic framework, a conceivable solution is the ad-hoc replacement of the δ function by a suitable smearing function $f(r)$ which converges weakly towards the δ function in the appropriate limit, i.e., the use of a representation of the δ function. This procedure should be independent of the particular choice of the representation of the δ function. The model of Ref. [62] replaces the δ function by the smearing function

$$f(r) = \frac{1}{4\pi r_0^2} \frac{\exp(-r/r_0)}{r} ,$$

satisfying

$$\lim_{r_0 \rightarrow 0} f(r) = \delta^{(3)}(\mathbf{r}) .$$

The dependence of the smearing-range parameter introduced here, r_0 , on the masses of the bound-state constituents is encoded in a purely phenomenological relationship [62]. The numerical values of the parameters of this model have been determined in Ref. [62] from a fit of the mesons known at that time, yielding, for the constituent quark masses, $m_u = m_d = 336$ MeV, $m_s = 575$ MeV, $m_c = 1845$ MeV, $m_b = 5235$ MeV, and, for the couplings in the potential, $\alpha_s = 0.31$, $a = 0.15$ GeV², $b = 0.956$ GeV, $c = 2.05$ GeV^{−1}.

For this quark interaction, the mass eigenvalues and eigenstates of the Breit–Fermi Hamiltonian are computed to high accuracy with the help of the numerical integration procedure for the solution of the Schrödinger equation developed in Ref. [65]. However, before applying the above quark model to the main target of the present investigation, viz., the system of charmed strange (D_s) mesons, we would like to confront this model’s predictions for the entire mass spectrum of the experimentally well-established mesons with results obtained in some representative different approaches. For this comparison, we choose the relativized quark potential model of Godfrey, Isgur, and Kokoski [15,16], and a relativistic quark model for mesons developed over the last decade by some Bonn group within the Lorentz-covariant framework of the Bethe–Salpeter equation [66–71].

In order to get some idea of the significance of our approach, that is, to estimate the accuracy of our predictions, and to compare them with the findings of different models, we introduce, as a measure of quality, a quantity Q , defined as the average of the square of the relative error of the theoretical result $M_i^{(\text{th})}$ for the mass of the composite particle $i = 1, 2, 3, \dots$ with respect to the corresponding experimentally measured value $M_i^{(\text{exp})}$:

$$Q \equiv \frac{1}{N} \sum_{i=1}^N \left(\frac{M_i^{(\text{th})} - M_i^{(\text{exp})}}{M_i^{(\text{exp})}} \right)^2 .$$

The smaller the numerical Q value, the higher is the accuracy of the adopted approach. For the numerical comparison of the models, we take into account their predictions for $N = 24$ energy levels corresponding to the well-established mesons π , $\rho(770)$, $\phi(1020)$, $\pi(1300)$, $\rho(1450)$, K , $K^*(892)$, D , $D^*(2010)$, D_s , D_s^* , B , B^* , B_s , $\eta_c(1S)$, $J/\psi(1S)$, $\eta_c(2S)$, $\psi(2S)$, $\psi(4040)$, $\psi(4415)$, $\Upsilon(1S)$, $\Upsilon(2S)$, $\Upsilon(3S)$, and $\Upsilon(4S)$. Numerically we find for the quality measure Q in our approach $Q = 0.019\%$, for the Godfrey–Isgur–Kokoski model [15,16] $Q = 0.041\%$, which is roughly twice as large as our Q value, and from the results presented in Refs. [68,70] for the Bethe–Salpeter treatment $Q = 0.28\%$, which is more than an order of magnitude larger than our Q value. From these numbers, we feel entitled to conclude that this quark potential model is definitely competitive (at least).

4 Results, Discussion, and Conclusions

Table 2 gives the predictions of the quark model sketched in Sec. 3 for the masses of the lowest-lying six ($c\bar{s}$) bound states expected by the spectroscopic classification of Sec. 2. Notice that we still adopt the model parameter values determined in the fit of Ref. [62]. The formalism of the Breit–Fermi Hamiltonian is more closely related to L - S coupling. Consequently, the states are labelled here by the usual spectroscopic notation $n^{2S+1}\ell_J$, where n denotes the principal quantum number, S is the quantum number of the total spin $\mathbf{S} = \mathbf{s}_Q + \mathbf{s}_{\bar{q}}$, and S, P, D, \dots indicate orbital angular momentum $\ell = 0, 1, 2, \dots$. There are two $J^P = 1^+$ states: the spin-singlet 1P_1 state and the spin-triplet 3P_1 state. Because of non-diagonal contributions of the relativistic corrections in the Breit–Fermi Hamiltonian, the two $J^P = 1^+$ states do not constitute mass eigenstates. The physical states associated to the lower (l) or the higher (h) mass eigenvalue, respectively, of the Breit–Fermi Hamiltonian are called $D_{s1}^{(l)}$ and $D_{s1}^{(h)}$. The latter states are obtained from the states 1P_1 and 3P_1 by an orthogonal transformation involving some mixing angle θ :

$$\begin{aligned} |D_{s1}^{(l)}\rangle &= \cos\theta |^1P_1\rangle - \sin\theta |^3P_1\rangle, \\ |D_{s1}^{(h)}\rangle &= \sin\theta |^1P_1\rangle + \cos\theta |^3P_1\rangle. \end{aligned}$$

In our quark model we find for this 1P_1 - 3P_1 mixing angle the numerical value $\theta = 44.7^\circ$.

Moreover, Table 2 compares the predictions of our approach with, on the one hand, the experimentally observed masses [14] of all previously known D_s mesons and, on the other hand, the results of the four (nonrelativistic) potential models [15–19] explicitly mentioned in the Introduction and of the Bonn Bethe–Salpeter model [66–71]. First of all, we note that all predictions of the Lorentz-covariant Bethe–Salpeter framework do not substantially differ from the potential-model ones; this indicates that the problems with the new D_s states are not just a consequence of the nonrelativistic approximation. Our results show a satisfactory agreement with experiment and are within the range of predictions of the other models. We are forced to conclude that the model of Ref. [62], too, cannot explain the masses of the new D_s resonances; this discrepancy becomes still more distinct when inspecting the *binding energy* of the quark–antiquark bound states.

We may get a clue to the correct identification of the heavier of the two new states, the $D_{sJ}(2463)^\pm$ meson, by inspection of the center of gravity (COG) of the $\ell = 1$ states. For P-wave states, the mass $M(\text{COG})_P$ of the center of gravity is defined according to

$$M(\text{COG})_P \equiv \frac{1}{9} [M(^3P_0) + 3 M(^3P_1) + 5 M(^3P_2)].$$

Table 2: Energy levels (in units of MeV) of the lowest S- and P-wave states, identified by the usual spectroscopic notation $n^2S^{+1}\ell_J$, of the charmed, strange quark–antiquark (D_s meson) system predicted by the nonrelativistic quark model developed in Ref. [62]; the two mass eigenstates of the two-state $J^P = 1^+$ P-wave subsystem corresponding to lower and higher mass, formed by linear combinations of the spin-singlet 1P_1 state and the spin-triplet 3P_1 state, are denoted by $D_{s1}^{(l)}$ and $D_{s1}^{(h)}$, respectively. For comparison, the experimental values and some results of typical quark models [16–19] are listed too. Since Refs. [16] and [19] round their results to 10 MeV zeros are added where necessary. The $D_{s1}(2536)^\pm$ of mass (2535.35 ± 0.6) MeV is one of the $J^P = 1^+$ states $D_{s1}^{(l)}$ and $D_{s1}^{(h)}$.

State	Experiment	This Work	Ref. [16]	Ref. [18]	Ref. [19]	Ref. [17]	Ref. [70]
1^1S_0	1968.1 ± 0.5	1963	1980	1968.8	1940	1965	1969
1^3S_1	2111.9 ± 0.7	2099	2130	2110.5	2130	2113	2116
1^3P_0		2446	2484	2387.8	2380	2487	2464
$D_{s1}^{(l)}$		2515	2550	2521.2	2510	2535	2506
$D_{s1}^{(h)}$		2527	2560	2536.5	2520	2605	2506
1^3P_2	2572.4 ± 1.5	2561	2590	2573.1	2580	2581	2552

Within the framework of (nonrelativistic) interaction-potential models for two-particle bound states, a *perturbative* inclusion of the spin–spin interaction implies the equality of the P-wave center-of-gravity mass $M(\text{COG})_P$ and the mass $M(^1P_1)$ of the 1P_1 state:

$$M(\text{COG})_P = M(^1P_1) \quad (\text{perturbatively}) .$$

In contrast to this, in a *nonperturbative* treatment of the spin–spin interaction by, e.g., some smearing of Dirac’s δ function—which should come closer to reality—the P-wave center-of-gravity mass $M(\text{COG})_P$ will be larger than the mass $M(^1P_1)$ of the 1P_1 state:

$$M(\text{COG})_P > M(^1P_1) \quad (\text{nonperturbatively}) .$$

Let’s neglect the mixing of the two $J^P = 1^+$ states necessary to get the physical states.

- Identifying the 3P_1 state with the $D_{s1}(2536)^\pm$ of mass $(2535.35 \pm 0.34 \pm 0.5)$ MeV, and the 1P_1 state with the $D_{sJ}(2463)^\pm$ meson of mass (2459.0 ± 1.1) MeV, entails

$$M(\text{COG})_P^{(\text{exp})} = 2531.7 \text{ MeV} > M(^1P_1) = M(D_{sJ}(2463)^\pm) = 2459.0 \text{ MeV} ,$$

in accordance with the above general theoretical expectations for this inequality.

- Identifying the 3P_1 state with the $D_{sJ}(2463)^\pm$ meson of mass (2459.0 ± 1.1) MeV, and the 1P_1 state with the $D_{s1}(2536)^\pm$ of mass $(2535.35 \pm 0.34 \pm 0.5)$ MeV, entails

$$M(\text{COG})_P^{(\text{exp})} = 2506.3 \text{ MeV} < M(^1P_1) = M(D_{s1}(2536)^\pm) = 2535.35 \text{ MeV} ,$$

in contradiction to the above general theoretical expectation for this inequality.

This simple observation may be regarded as a hint that the new $D_{sJ}(2463)^\pm$ resonance is predominantly a 1P_1 state.

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